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Impact of Higher Order and Soft Gluon Corrections on the Extraction of Higher Twist Effects in DIS

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Abstract

The impact of recently calculated next-to-next-to-leading order QCD corrections and soft gluon resummations on the extraction of higher twist contributions to the deep-inelastic structure function F_2 is studied using the BCDMS and SLAC data. It is demonstrated to which extent the need for higher twist terms is diminishing due to these higher order effects in the kinematical region, $0.35 \leq x \leq 0.85$ and $Q^2 > 1.2 \text{ GeV}^2$, investigated. In addition, theoretical uncertainties in the extraction of higher twist contributions are discussed, and comparisons to results obtained previously are made.

Deep-inelastic lepton-hadron scattering (DIS) data are an integral part of all global QCD analyses of parton distribution functions [1, 2]. To guarantee that only the leading-twist part of the parton densities is extracted usually only data with sufficiently high momentum transfer Q^2 and invariant mass $W^2 = Q^2(1/x - 1)$ are selected, e.g., $Q^2 > 2 \text{ GeV}^2$ and $W^2 > 10 \text{ GeV}^2$ [1]. These universal distributions can then be applied to other hard processes by virtue of the factorization theorem. Contributions of higher twist (HT) are expected to become increasingly important as $x \rightarrow 1$ as shown, e.g., in the infrared renormalon approach [3], but they are suppressed by additional powers of $1/Q^2$ with respect to the logarithmic Q^2 dependence of the leading twist contribution. Hence HT operators can be safely neglected in all conventional global QCD analyses [1, 2] due to the aforementioned cuts on Q^2 and, in particular, on W^2 .

However, twist-four contributions are interesting in themselves and can provide valuable insight into higher quark-quark and quark-gluon correlators inside the nucleon. Not very much is presently known about these correlators apart from some studies in the framework of lattice QCD [4]. Recent measurements of the transversely polarized DIS structure function g_2 [5] may indicate that at least some HT are small¹ for an averaged Q^2 of about $2 \div 3 \text{ GeV}^2$ even at fairly large values of x . A better understanding of HT contributions, in particular their importance in describing low Q^2 DIS data, is important in many respects. A wealth of low Q^2 DIS data presently discarded in global QCD analyses would open up. Secondly, future experiments, e.g., the ‘CEBAF @ 12 GeV’ program, focus on measurements at high x at comparatively low Q^2 where HT are expected to be relevant. Finally, QCD analyses of longitudinally polarized DIS data would benefit. Here the available data are too scarce to allow for sufficiently ‘safe’ cuts on Q^2 and W^2 when extracting leading twist polarized parton densities, and the importance of HT in relating the measured spin asymmetry $A_1 \simeq g_1/F_1$ (where HT may partly cancel in the ratio) to the structure function g_1 is still an open issue, see, e.g. Ref. [6].

At present, data on the unpolarized structure function F_2 obtained quite some time ago by the CERN-BCDMS and various SLAC experiments [7] still offer the best testing

¹Measurements of g_2 provide a unique possibility to disentangle its twist-three part since the leading twist part can be entirely expressed in terms of the rather well known structure function g_1 .

ground for HT studies since they cover a wide kinematical range up to large x and down to low values of Q^2 with sufficiently good statistical accuracy. Therefore several attempts have been already made to disentangle leading and higher twist contributions to F_2 [8-14]. As anticipated, HT effects become increasingly important as $x \rightarrow 1$. Based on a partial next-to-next-to-leading order (NNLO) QCD analysis it was argued [13] that the need for HT is diminishing greatly when going from the NLO to the NNLO of QCD². A first complete NNLO analysis [14] could not confirm this observation. However, in that paper no distinction between different sources of HT terms was made, and it could be that most of the remaining HT in NNLO is of purely kinematical origin and hence calculable (see discussion below Fig. 1).

In the remainder of the paper we further extend these studies making use of the NNLO coefficient functions [17] and recent estimates for the full x dependence of the NNLO splitting functions [18] based on the integer Mellin n moments calculated in [19]. Another important ingredient of our analyses will be soft gluon resummations (SG) which have not been considered so far in extractions of HT contributions to DIS. The quark coefficient functions $C_{1,2}^q$ in DIS which link the quark distributions to the structure functions $F_{1,2}$ exhibit a large x enhancement of the form $[\ln^{l-1}(1-x)/(1-x)]_+$ where $l = 1, \dots, 2k$ in the $\mathcal{O}(\alpha_s^k)$ approximation to $C_{1,2}^q$ which needs to be resummed to all orders [20, 21]. Finally, we compare the outcome of our analyses with results already available in the literature. It should be stressed that we mainly focus on a qualitative comparison of HT contributions to F_2 extracted in NLO, NNLO, and NNLO including SG and do not attempt to provide a full set of parton densities or to extract $\alpha_s(M_Z)$ from BCDMS and SLAC F_2 data as was done, e.g., in [8, 10]. Instead we study the various sources of theoretical uncertainties in the determination of HT contributions to DIS.

The DIS structure function F_2 can be expanded in $1/Q^2$ as

$$F_2(x, Q^2) = F_2^{(2)}(x, Q^2) + F_2^{(4)}(x, Q^2)/Q^2 + \mathcal{O}(1/Q^4) \quad (1)$$

where $F_2^{(t)}$ denotes the contribution of twist- t . The leading twist part $F_2^{(2)}$ ‘factorizes’

²It should be noted that similar studies have been made also for the charged current structure function xF_3 [15, 16]. In [16] it was also observed that NNLO contributions largely remove the need for HT at large x .

into a convolution of a perturbatively calculable coefficient function C_2 and some non-perturbative parton density combination f which can only be determined from experiment so far:

$$F_2^{(2)}(x, Q^2) = x[C_2(\alpha_s(\mu_r^2), \frac{Q^2}{\mu_f^2}, \frac{\mu_f^2}{\mu_r^2}) \otimes f(\mu_f^2, \mu_r^2)](x) \quad . \quad (2)$$

Once the non-perturbative input f is fixed at some reference scale μ_0 its μ_f dependence is fully predicted by the well-known DGLAP evolution equation which schematically reads

$$\frac{d}{d \ln \mu_f^2} f(x, \mu_f^2, \mu_r^2) = [\mathcal{P}(\alpha_s(\mu_r^2), \frac{\mu_f^2}{\mu_r^2}) \otimes f(\mu_f^2, \mu_r^2)](x) \quad , \quad (3)$$

where μ_f and μ_r denote the factorization and renormalization scales, respectively. For simplicity we have limited ourselves in Eqs. (2) and (3) to the non-singlet sector which is all what is needed for our analyses as will be discussed below.

The coefficient function C_2 in (2) is known up to NNLO [17] while only the first integer moments of the evolution kernels \mathcal{P} in three-loop order have been calculated so far [19]. However, based on these results and further constraints on \mathcal{P} estimates for the full x dependence of the NNLO kernels have been derived recently [18]. The residual uncertainties on the kernels were shown to be extremely small [18] in the large x region, $x \gtrsim 0.3$, we are interested in. The relevant coefficients of the QCD beta function up to β_2 which govern the running of $\alpha_s(\mu_r)$ are given in [22]. The appropriate matching conditions at flavor thresholds can be found in [23]. We are using the $\overline{\text{MS}}$ scheme throughout this paper.

There are two sources of power corrections in $1/Q^2$ contributing to $F_2(x, Q^2)$ in (1) beyond twist-2. The first one is of purely kinematical origin and can be entirely expressed in terms of the leading twist contribution $F_2^{(2)}(x, Q^2)$ in (2). It only takes into account effects of the so far neglected non-zero target mass M and approximately behaves like $x^2 M^2/Q^2$. Hence it gives a sizable contribution to F_2 at large x whenever Q^2 is of $\mathcal{O}(M^2)$ or smaller. The ‘target mass corrected’ (TM) expression for $F_2^{(2)}(x, Q^2)$ is known for a long time and reads [24]

$$F_2^{\text{TM}}(x, Q^2) = \frac{x^2/\xi^2}{r^{3/2}} F_2^{(2)}(\xi, Q^2) + 6 \frac{M^2}{Q^2} \frac{x^3}{r^2} \int_{\xi}^1 \frac{d\xi'}{\xi'^2} F_2^{(2)}(\xi', Q^2)$$

$$+12\frac{M^4}{Q^4}\frac{x^4}{r^{5/2}}\int_{\xi}^1 d\xi' \int_{\xi'}^1 \frac{d\xi''}{\xi''^2} F_2^{(2)}(\xi'', Q^2) \quad , \quad (4)$$

with $\xi = 2x/(1 + \sqrt{r})$ and $r = 1 + 4x^2 M^2/Q^2$. It is sometimes preferred to use F_2^{TM} only after expanding (4) in powers of M^2/Q^2 and retaining only the leading term. We will use both, the full and the expanded expressions for F_2^{TM} in our fits. Any differences somehow represent part of the theoretical uncertainty in the extraction of higher twists from DIS data.

The other source of power corrections cannot be related to $F_2^{(2)}(x, Q^2)$ in general and provides important new insight into the QCD dynamics of higher quark-quark and quark-gluon correlations in the nucleon about which almost nothing is known yet. Also the Q^2 evolution of these twist four operators has not been calculated yet³. Therefore one has to fully rely on some model here. One interesting possibility is the infrared renormalon approach [3] which allows to calculate explicitly the power corrections to the coefficient function C_2 in (2), i.e., to predict the x shape of $F_2^{(4)}(x, Q^2)$ up to some unknown normalization factor, which has to be determined from data, see [25] for details. This approach was used in some of the previous analyses of higher twists in DIS [10, 12, 13]. It turned out that, although the renormalon model could reproduce the general trend that HT increase as $x \rightarrow 1$, the predicted x shape was somewhat off (see, e.g., Fig. 2 in [10]). Therefore we prefer not to fix the x shape of $F_2^{(4)}(x, Q^2)$ and instead use a simple multiplicative ansatz⁴ as was also employed in the first HT analysis by Milsztajn and Virchaux [8] as well as in [12, 13]. For each x bin of the data we simply introduce a factor C_{HT} which is independent of Q^2 ,

$$F_2(x, Q^2) = F_2^{\text{TM}}(x, Q^2) \left[1 + \frac{C_{\text{HT}}(x)}{Q^2} \right] \quad , \quad (5)$$

i.e., $F_2^{(4)}(x, Q^2)$ is *approximated* to have the same logarithmic dependence on Q^2 as $F_2^{(2)}(x, Q^2)$.

Finally, we also study the influence of soft gluon resummations for the coefficient func-

³ Even if the evolution of twist four operators would become available, it would be extremely difficult to make use of it because the twist-4 operators which are accessible in DIS mix under evolution with other twist-4 operators which do not contribute to DIS and have to be taken from elsewhere.

⁴Sometimes $F_2^{(4)}(x, Q^2)$ is defined in such a way that it does not exhibit any Q^2 dependence, i.e., $F_2(x, Q^2) = F_2^{\text{TM}}(x, Q^2) + \tilde{C}_{\text{HT}}(x)/Q^2$ [11].

tion C_2 in (2) [20, 21] on the extraction of HT from DIS data. Since these resummations are operative also in the large x region where HT become important they may have a sizable impact on the size of the extracted HT coefficients C_{HT} in (5). The SG take care of potentially large logarithms of the form $[\ln^{l-1}(1-x)/(1-x)]_+$ where $l = 1, \dots, 2k$ in the $\mathcal{O}(\alpha_s^k)$ approximation to C_2 which threaten to spoil the convergence of the perturbative expansion by resumming them to all orders in α_s . These resummations have been recently pushed up to next-to-next-to-leading logarithmic terms [21] and can be straightforwardly implemented in the Mellin n moment space which we also use for solving the evolution equations (3).

Before we turn to our numerical results let us specify the ansatz for the parton density function f in (2) and the data sets and cuts used in our analyses. Since we are only interested in the impact of NLO, NNLO, and SG corrections on the size of the higher twist coefficients C_{HT} in (5) we can limit ourselves to the large x region by taking into account only the BCDMS and SLAC data (using the BCDMS x binning) with $x \geq 0.35$ [7]. This is also the region where the estimates for the NNLO kernels \mathcal{P} work extremely well [18]. As in [10-13] we then apply the non-singlet (‘valence’) approximation for F_2 and only a single combination of parton densities f is required for proton target data⁵ (we will comment below on a possible ‘contamination’ of the scaling violations of F_2 at large x due to singlet and gluon contributions). In addition we select only data with $Q^2 \geq 1.2 \text{ GeV}^2$ to stay away from the resonance region and the region where the perturbative expansion is expected to break down since α_s becomes too large. Six bins in x remain after our cuts, so we have to determine six different HT parameters $C_{\text{HT}}(x = 0.35), \dots, C_{\text{HT}}(x = 0.85)$ in (5) from the fit to the data.

We parametrize the input distribution $f(x, \mu_0^2)$ in (2) by the standard ansatz

$$f(x, \mu_0^2) = Nx^\alpha(1-x)^\beta(1 + \gamma_1\sqrt{x} + \gamma_2x) \quad (6)$$

using $\mu_0 = 1 \text{ GeV}$ as the initial scale. It turns out, however, that all fits are rather insensitive to γ_1 and γ_2 , so we will neglect these terms in the following. $\alpha_s(\mu_r)$ is always

⁵We refrain from using deuterium data [7] as well. They are equally precise than proton data but may suffer from nuclear binding effects. We also discard SLAC data with $x > 0.85$ [26] which are close to or in the resonance region.

computed by solving its renormalization group equation exactly in NLO or NNLO since the approximate formula [27] is not sufficient for small scales. For simplicity we always choose $\mu_r = \mu_f = Q$. The evolution equations (3) are most conveniently solved in Mellin n space. There are, however, different ways how to actually truncate the solution at a given order k in perturbation theory ($k = 1$: NLO, $k = 2$: NNLO), see, e.g., [28] for a detailed discussion. On the one hand one can have a sort of iterative solution where all orders in α_s still contribute, on the other hand one can impose strict power counting by keeping only terms up to $\mathcal{O}(\alpha_s^k)$. Both approaches are widely in use in global QCD analyses and differ, of course, only by terms of $\mathcal{O}(\alpha_s^{(k+1)})$. This ‘freedom’ also represents part of the theoretical uncertainty and we will perform our fits using different solutions to (3).

We should also mention that we have always added the statistical and systematic errors in quadrature as is commonly done in most QCD analyses of parton densities. In [8] *some* of the systematic errors of the BCDMS data have been combined to a so called ‘main systematic error’ [29]. The BCDMS data points are then allowed to float within this error with a common normalization factor to be determined from the χ^2 fit [8]. The BCDMS data shifted in this way are used in some of the global QCD analyses. Unless stated otherwise, we refrain from using these data [8] as they are strongly biased by the theoretical input that went into the analysis, and we stick to the original data sets [7] instead. In particular, it should be noticed that in [8] data with Q^2 values as low as 0.5 GeV^2 were taken into account and that only the approximate solution for the running of α_s [27] was used which is known to be way off the exact solution at such small scales. In addition, it has been shown recently in [10] that the introduction of the main systematic error is not a rigorous method to deal with systematic errors and that correlations should be fully taken into account to obtain reliable errors for the parameters extracted from a χ^2 analysis. Nevertheless, as we are mainly interested in the qualitative effects of higher order corrections we treat all errors as uncorrelated which greatly simplifies the analysis and should not effect the general features of our results. The statistical meaning of the χ^2 function is totally obscured in QCD analyses of parton densities anyway, and errors are very difficult to quantify [30]. Therefore the errors for the C_{HT} as obtained in our χ^2 fits

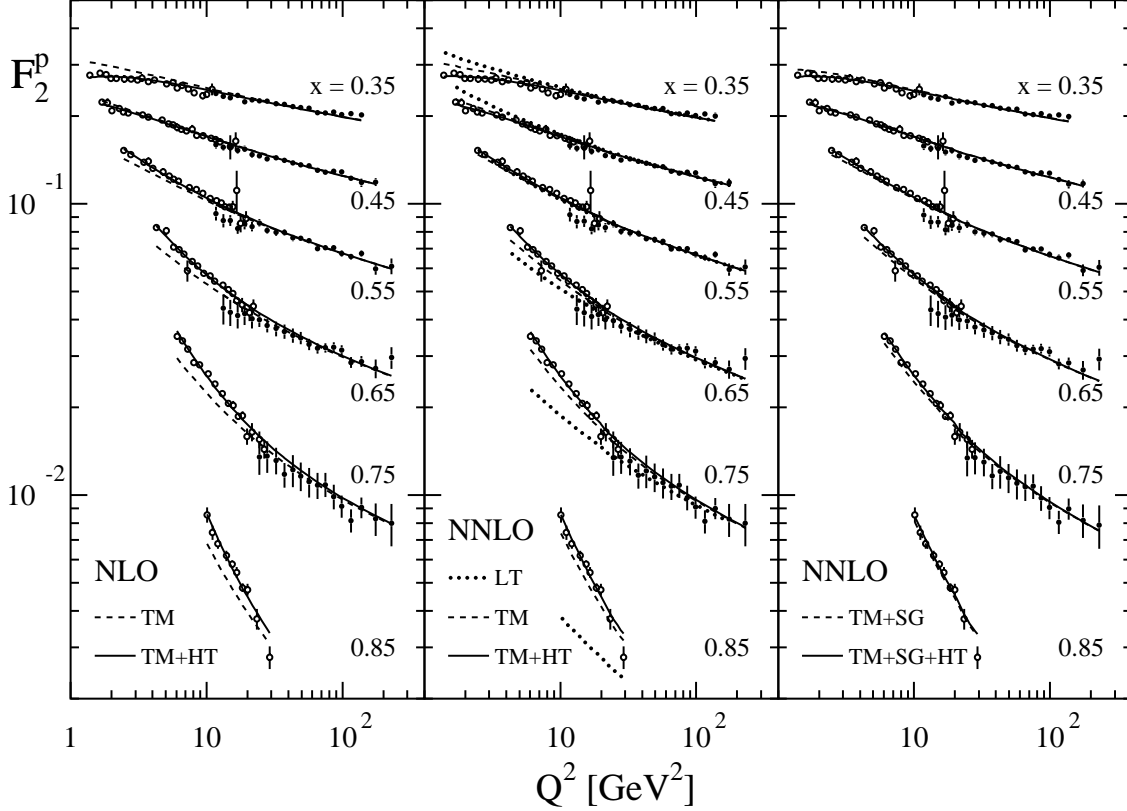


Figure 1: Comparison of the results of our χ^2 fits (solid lines) in NLO and NNLO of QCD and in NNLO including SG with the SLAC and BCDMS proton target data [7] (open and full circles, respectively) using $\alpha_s(M_Z) = 0.117$. The dashed lines refer to the target mass corrected part $F_2^{\text{TM}}(x, Q^2)$ only with the additional higher twist factor $(1 + C_{\text{HT}}(x)/Q^2)$ in Eq. (5) being omitted. The leading twist (LT) part $F_2^{(2)}$ without TM is shown as well in the center part (dotted lines) for illustration.

should not be taken too literally or be regarded as 1σ -errors. As in [8, 12, 13] we allowed for a global normalization shift of the BCDMS data with respect to the SLAC data.

In Fig. 1 we compare the results of our χ^2 analyses of F_2 according to Eq. (5) (solid lines) with the SLAC and BCDMS proton target data [7] using $\alpha_s(M_Z) = 0.117$. Very good agreement with very similar values of χ^2 is achieved for the fits in NLO and NNLO QCD as well as for a fit in NNLO including the SG mentioned above. The size of the required dynamical HT contributions can be inferred from comparing the dashed lines, which refer to the TM corrected part of F_2 only with $(1 + C_{\text{HT}}(x)/Q^2)$ in Eq. (5) being

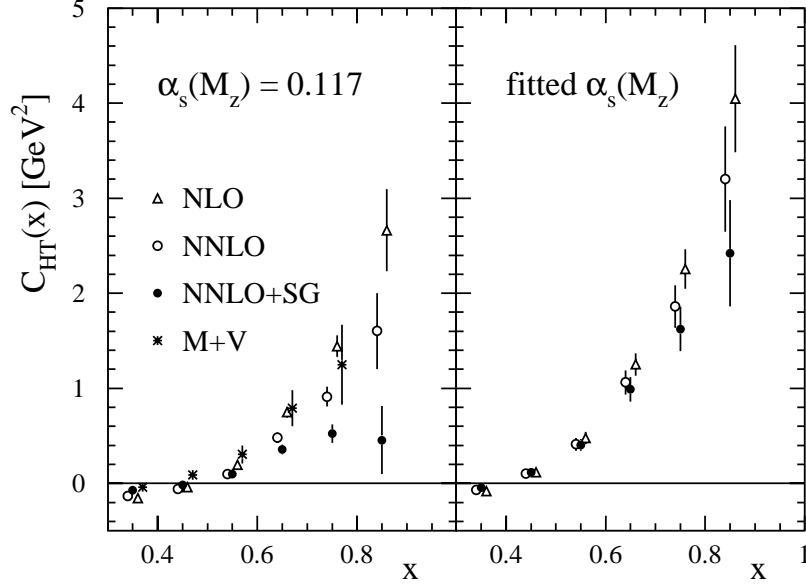


Figure 2: Extracted higher twist parameters C_{HT} in six x -bins according to Eq. (5) for fixed and fitted values of $\alpha_s(M_Z)$. On the l.h.s. also the results obtained in [8] are shown for comparison (crosses). For each x -bin the results are slightly horizontally displaced to avoid overlapping error bars.

omitted, with the full results (solid lines). Clearly the need for HT contributions is diminishing when going from NLO to NNLO. A further reduction, most noticeable at $x = 0.85$, is observed when the SG are taken into account as well. In the center part of Fig. 1 we also give the leading twist (LT) result $F_2^{(2)}$ without target mass correction. A comparison with the corresponding TM and full results shows that a major part of the HT contributions required to describe the data is of purely kinematical rather than dynamical origin and hence calculable. The extracted higher twist coefficients C_{HT} for the six x bins are shown separately on the left hand side (l.h.s.) of Fig. 2.

It should be mentioned that we have also tried to obtain $\alpha_s(M_Z)$ from the fits rather than setting it to a fixed value. Unfortunately, the values we obtained turned out to be unacceptably small, $\alpha_s(M_Z) \approx 0.105$, way outside the current world average even within the error bars. The gross feature that the HT parameters get smaller when going to higher orders of perturbation theory is common to all results but less pronounced for

$\alpha_s(M_Z) \approx 0.105$ as can be seen by comparing the left and the right hand sides of Fig. 2. This also demonstrates the strong correlation between the size of the C_{HT} coefficients and the value of $\alpha_s(M_Z)$: the smaller the $\alpha_s(M_Z)$ the larger the HT contribution as was already pointed out in [10].

We have investigated the reason for not obtaining a more common value for $\alpha_s(M_Z)$ from our fits in particular because reasonable values for $\alpha_s(M_Z)$ have been obtained in previous analyses [8, 10]. The latter fits differ, however, in various aspects from our ones. In [8] a NLO fit to all available BCDMS and SLAC proton and deuterium data was performed⁶ in the range $Q^2 > 0.5 \text{ GeV}^2$ and $0.07 \leq x \leq 0.75$. If we extend the range of data included in our fit to $0.07 \leq x$, but still requiring $Q^2 > 1.2 \text{ GeV}^2$, and introduce also a singlet and gluon contribution to F_2 , we obtain much more reasonable values for $\alpha_s(M_Z)$ in NLO as well. This indicates that α_s cannot be extracted from large x data alone and/or that singlet and gluon distributions are still important when studying scaling violations of F_2 for $x > 0.35$, i.e., that a pure non-singlet approximation for F_2 is not sufficient here. Another possible explanation for the small value of $\alpha_s(M_Z)$ may come from our simplified treatment of the systematical errors. To investigate this further we have performed similar fits using the data in [8] which include the BCDMS main systematic error shift mentioned above. In NLO this also leads to a significantly increased value for $\alpha_s(M_Z)$ but almost no changes occur in NNLO. This does not come as a surprise since the main systematic error shift in [8] was obtained only from a NLO fit, and it strongly depends on the theoretical input. It seems to be reasonable to assume that an $\alpha_s(M_Z)$ of about $0.116 \div 0.118$ would be obtained in our fits once the correlations between different sources of systematical errors are properly taken into account as in [10, 11]. This is however far beyond the scope of our more qualitative studies. Data at smaller values of x may also be important but a reliable NNLO analysis can only be performed once the complete NNLO kernels become available. As in [12, 13], we therefore prefer to fix $\alpha_s(M_Z)$ in all our fits to 0.117 rather than taking the fitted value which just seems to be an artifact of the approximations used

⁶ As mentioned above, the approximate expression for $\alpha_s(\mu_r)$ [27] was used in [8] which is not valid for small scales. In addition not all details of their analysis are specified in [8]. Therefore it is difficult to actually compare the result for $\alpha_s(M_Z)$ (and the HT coefficients) obtained in [8] with our NLO results on a quantitative level.

in our analyses.

Before we turn to a comparison with results obtained previously in the literature, let us briefly focus on the theoretical uncertainties in the extraction of C_{HT} [and $\alpha_s(M_Z)$] from DIS data. We have already mentioned that there are different ways of how to treat the TM in (4) and that there is no unique way of actually solving the evolution equations (3) to a given order. We have performed several fits to take this ambiguity into account and it turns out that the gross features of $C_{\text{HT}}(x)$ are basically unchanged but the precise values of the $C_{\text{HT}}(x)$ do depend to a certain extent on the details of the analysis. For simplicity we have always chosen $\mu_r = \mu_f = Q$ but nothing prevents us from varying these two scale independently around Q which would also alter the values of the $C_{\text{HT}}(x)$. The dependence on the factorization scheme and on μ_f can be removed if one expresses the scaling violations of F_2 in terms of F_2 itself using ‘physical’ evolution kernels, see, e.g., Ref. [31]. Another major uncertainty which is difficult to quantify stems from the form of the chosen ansatz (5) for the HT terms (see the discussion above). In addition we cannot take into account the correct Q^2 dependence of the HT terms since the relevant anomalous dimensions are not known (cf. footnote 3). Apart from studies of the scale dependence, see, e.g., Refs. [8, 11], none of the other sources of theoretical uncertainties have been included in any of the combined extractions of HT terms and $\alpha_s(M_Z)$ from DIS data. Since $C_{\text{HT}}(x)$ and $\alpha_s(M_Z)$ are highly correlated we therefore argue that errors on $\alpha_s(M_Z)$ from those types of analyses are perhaps seriously underestimated and of limited use (in particular when one takes into account also the obscured statistical significance of χ^2 in a QCD analysis [30]). To determine $\alpha_s(M_Z)$ from DIS it seems to be much safer to introduce first appropriate cuts on Q^2 and W^2 to remove the kinematical region which is contaminated by HT and then use the obtained $\alpha_s(M_Z)$ value (or that from a global QCD analysis) to determine the $C_{\text{HT}}(x)$ afterwards by relaxing the cuts and keeping $\alpha_s(M_Z)$ fixed.

Finally, let us compare our results with previous ones available in the literature. On the l.h.s. of Fig. 2 we also show the HT coefficients obtained in the NLO fit in [8]. Despite the differences in both analyses (see also footnote 6) reasonable agreement is achieved

with our NLO results for C_{HT} . We cannot directly compare with the C_{HT} in [10, 11] since their ansatz differs from (5), see also footnote 4, but the main features are the same, in particular the change of sign in C_{HT} at around $x \simeq 0.45$. As already mentioned in the introduction, it was argued in [13] that NNLO corrections reduce the need for HT contributions to F_2 . However, in [13] only the NNLO coefficient functions were used and the parton densities were taken from a NLO global analysis without refitting them. Each x bin was assigned a separate normalization factor, called ‘floating factor’, to substitute for the necessary refit of the parton densities [13], and the HT coefficients were taken from the IR renormalon model [25]. Nevertheless, we confirm their conclusions. In addition, SG seem to further reduce the need for HT at large x .

In Fig. 3 our fitted input distributions f are compared with the corresponding combination of parton densities from a recent *global* NLO QCD analysis [1] which to a large extent excludes the kinematical region where HT become operative. Over the entire x range included in our analysis, $0.35 \leq x \leq 0.85$, our NLO f differs from the MRST result by less than 10%. Only at NNLO with SG included, our f turns out to be significantly smaller than the MRST result as $x \rightarrow 1$. This gives an indication of the importance of SG at large x . Since the resummations increase the coefficient functions smaller parton densities are needed in that region to describe the data.

To summarize, we have studied the impact of higher order corrections on the extraction of HT contributions to the DIS structure function F_2 . The size of the required HT terms is diminishing when NNLO corrections are taken into account as was recently argued in the literature. However, even at NNLO with SG still a sizable, positive HT term is needed at large x while for $x \lesssim 0.45$ a small negative HT contribution is required. SG resummations tend to further reduce the size of the HT contributions and should be therefore included in future analysis of upcoming data at large x , e.g., from ‘CEBAF @ 12 GeV’. It was argued that simultaneous extractions of $\alpha_s(M_Z)$ and HT terms in DIS may suffer from large theoretical uncertainties due to the strong model dependence and the unknown Q^2 evolution of the HT contribution.

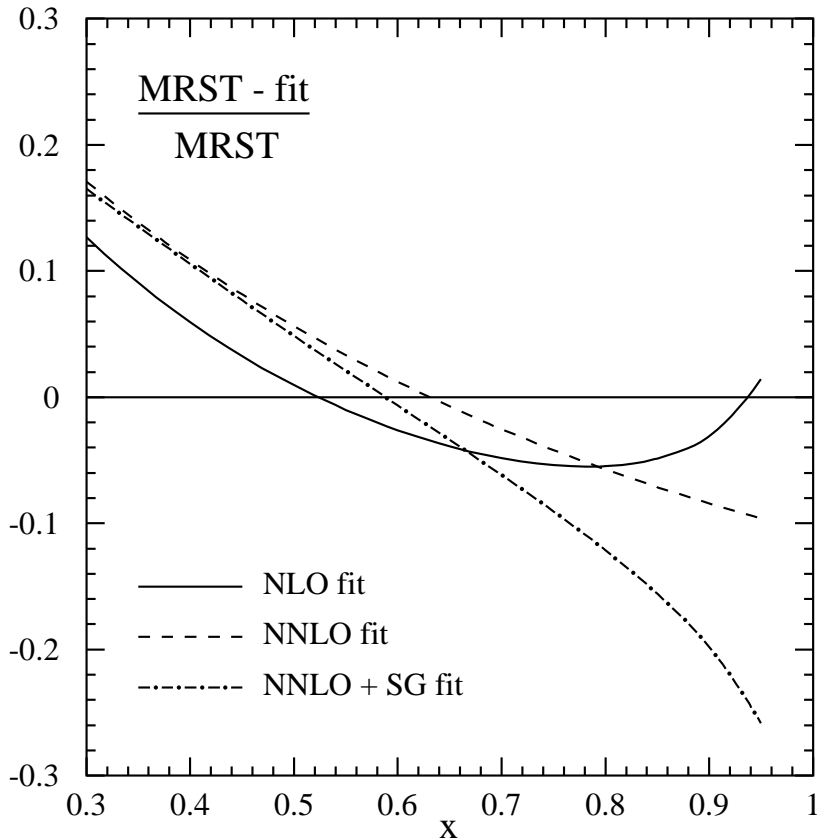


Figure 3: Comparison of our fitted input distribution f according to Eq. (6) with the corresponding non-singlet (valence) combination of a recent MRST global analysis [1]. Shown are the deviations from the MRST ‘central fit’ at the input scale $\mu_f = \mu_0 = 1$ GeV.

Acknowledgments

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